Spontaneous X-rays Exclude Gaussian Covariant Classical–Quantum Gravity

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Abstract

In classical–quantum gravity the metric stays classical and complete positivity ties the decoherence of quantum matter to the diffusion of the metric. The covariant theory of Oppenheim and collaborators is consistent with interferometric, LISA Pathfinder and gravitational-wave data. Its matter noise at keV frequencies, not evaluated before, conflicts with XENONnT's search for spontaneous X-rays, by two arguments that share the trade-off, the atom and the data but fail under different escapes. First, the covariant diffusion makes the conformal mode dynamical; near the light cone slow matter feels it through its rest mass, while Maxwell's equations do not. For translation-invariant kernels complete positivity holds momentum by momentum, and on a conserved near-null source it forces matter noise $s \geq \frac{2}{3}|f|/p^2$ for every diffusion weight that is Lorentz invariant in the scalar sector, including every regulator or running coupling that depends on p^2 alone. Newton's law, tested from 1 mm to 1 m, then gives $\int s\, dp^2 \geq 9.2$ at every X-ray energy, against the XENONnT limit $C\int s < 7.7 \times 10^{-8}$: an exclusion by $\geq 1.1 \times 10^8$, which needs Lorentz invariance of the scalar sector only up to boosts $\gamma = 1.1 \times 10^7$. A frame in the scalar weight evades it. Second, general relativity's response to a conserved near-null source is a pure tensor with a light-cone pole. The response-squared bound holds for every positive diffusion weight. It assumes general relativity's tensor propagator at keV momenta but no transfer of static tests, and the Coulomb field of an atom gives the tensor source an order-one share. Our estimate of this margin is 1.7×10^9, so frames in the weight, including a dynamical foliation, do not shield. A classical theory that survives both needs a frame or a non-perturbative mass above $(0.34\ \mu\text{m})^{-1}$ in the scalar sector together with a tensor sector that leaves general relativity below about 10 eV (for a $z = 3$ Lifshitz dispersion) or a matter coupling blind to the atom's stress, or a frame carried by the matter coupling, or non-Gaussian noise that delivers its keV emission in multi-photon bursts.

1 Introduction

Is gravity quantized? Entanglement witnesses aim to answer through the entanglement a gravitational interaction can create [1, 2]. Classical–quantum (CQ) gravity answers the other way: it keeps the metric classical and couples it to quantum matter by a linear, completely

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positive (CP) dynamics [3]. The price is fixed by positivity. A classical field that responds to quantum matter must diffuse, and the matter must decohere, with the product of the two bounded below: the decoherence–diffusion trade-off [4]. In the Newtonian limit the least decoherence allowed is the Diósi–Penrose rate [5, 6, 7], whose spontaneous X-ray emission is bounded by underground detectors [8, 9, 10].

A covariant version exists. Its diffusion kernel is quadratic in curvature, of DeWitt form $\alpha R_{\mu\nu}R^{\mu\nu} - \beta R^2$ [11]; it is power-counting renormalisable with a positive scalar sector [12], and its scalar–vector–tensor decomposition shows a dynamical, stochastic conformal mode beside the spin-2 mode [13]. Its authors find it viable against matter-wave coherence, LISA Pathfinder and gravitational-wave data [13]. They find that with scale-independent couplings the integrated energy density of the stochastic gravitational waves would rule the theory out, and that asymptotic freedom [12] makes the coupling run [13]. Oppenheim and Panella expected anomalous heating bounds to be the most stringent [14], but did not compute them. They also noted that the tensor trade-off decoheres the strain components of the stress tensor, for example through $T_{ij}T^{ij}$, so that it is unclear whether heating constrains the tensor coupling [14]. Diósi showed that the heating of Lorentz-invariant Markovian decoherence and diffusion diverges [17, 18].

This paper computes the keV matter noise of covariant CQ gravity and confronts it with the XENONnT spectrum [10]. The idea fits in one line. Lorentz invariance maps the laboratory band on which Newton's law is tested, invariant lengths from 1 mm to 1 m, onto a thin strip next to the light cone at every X-ray energy; the trade-off forces noise on that strip; slow matter feels it through its rest mass and radiates it as keV photons. The number is

$$\frac{\int s\,du}{\text{XENONnT limit}} \;\geq\; \frac{\frac{2}{3}\ln 10^6}{7.7\times 10^{-8}} \;=\; 1.2\times 10^8, \tag{1}$$

with s the matter noise in units of the Diósi–Penrose kernel and $u = p^2$ the invariant momentum squared (Section 2). After a bound-electron envelope we quote the exclusion as $\geq 1.1\times 10^8$. It holds for every DeWitt coupling, every regulator and every running coupling that depends on p^2 alone, and it needs Lorentz invariance of the scalar sector only up to boosts $\gamma = 1.1\times 10^7$.

The scalar exclusion has a precise escape: a frame acting on the scalar sector. One of the authors' completions, which puts a unit timelike vector into the scalar weight [13], survives it; so does a preferred foliation of the Hořava type [24, 25]. Section 5 closes that escape with a second argument. General relativity's response to a conserved near-null source is a pure tensor with a light-cone pole. The response-squared form of the trade-off holds for every positive diffusion weight, so no frame in the weight removes it. It assumes general relativity's tensor propagator at keV momenta, but it needs no transfer of static tests through Lorentz boosts. The atom's Coulomb field gives the tensor source an order-one share. Its margin, 1.7×10^9, is an estimate, because the tensor conversion uses the elastic structure factor of an isolated atom. The two exclusions share several steps, which Section 5.5 lists; they are independent in the sense that different escapes defeat each.

What is not new. The dynamical scalar is Oppenheim and Sajjad's; its positivity is from [12]. So is the static scalar floor, a factor 25/9 above the tensor channel's in their product of diffusion and decoherence; Remark 3.1(iii) gives its square root, 5/3, in our units [13]. That the tensor channel decoheres the transverse-traceless anisotropic stress is stated in [14] and [13]. The trade-off, and the remark that heating bounds on decoherence bound diffusion through it, are from [4].

The divergence of Lorentz-invariant Markovian heating is Diósi's [17, 18]. The X-ray route was named in [14], and X-ray bounds on Newtonian CQ gravity follow from Tilloy and Diósi's identification of its least-decoherence limit with the Diósi–Penrose model [6]. Gravitational-wave strain noise [19] and classical–quantum scattering [20] have been used against covariant CQ models by other routes. The conversion of noise into photons is Manti *et al.*'s [21]. What we add is the near-cone face of the conformal mode, the α-independent floor that survives every p^2-only regulator, the tensor floor that survives every positive weight, the observation that an atom's Coulomb field makes the tensor channel of order one at keV, and the quantitative XENONnT bounds. We found no earlier X-ray bound on covariant CQ gravity.

Section 2 sets up the Gaussian CQ dynamics and derives the trade-off at fixed momentum. Section 3 gives the scalar floor, Section 4 its conversion and the XENONnT exclusion, Section 5 the tensor exclusion, Section 6 the objections and the scope, and Section 7 what remains.

2 Gaussian classical–quantum dynamics and the trade-off

2.1 Conventions

Signature $(-,+,+,+)$; $c=1$ in formulas, SI in numbers. For a four-momentum $p=(\omega,\mathbf{k})$,

$$u \equiv p^2 = k^2 - \omega^2, \tag{2}$$

so $u>0$ is spacelike and the light cone is $u=0$. Matter noise acts on the rest-mass density $\rho_m = -T^{\mu}{}_{\mu}$ and enters a branch pair through the decoherence exponent $-\frac{1}{2}\int d^4p\,(2\pi)^{-4}|\Delta\rho_m(p)|^2 S(p)$. We quote it in units of the Diósi–Penrose kernel,

$$s \equiv \frac{\hbar S}{4\pi G}, \qquad s_{\mathrm{DP}} = \frac{1}{k^2} \text{ at } C=1, \qquad s = \frac{N}{4\pi\ell_{\mathrm{P}}^2}, \quad 4\pi\ell_{\mathrm{P}}^2 = \frac{4\pi\hbar G}{c^3} = 3.28\times10^{-69}\ \mathrm{m}^2, \tag{3}$$

with N (m^4) a kernel per unit ρ_m^2. The Newton form factor $f(u)$ is the classically carried response relative to Newton's, $|f|=1$ where Newton's law is tested. A margin is the lower bound on $\int s\,du$ divided by the XENONnT upper limit on $C\int s\,du$; a margin above one excludes.

2.2 The trade-off

In continuous (Gaussian) CQ dynamics [15, 4, 11] a classical record A couples to quantum matter branches $m^{\pm}$ through three kernels: the decoherence kernel D_0, the diffusion D_2 of the record, and the drift D_1 that carries the mean field. The weight on a branch pair is

$$W = \exp\Big\{-\tfrac{1}{2}\Delta m\, D_0\, \Delta m - \tfrac{1}{2}(A - D_1\bar m)\, D_2^{-1}(A - D_1\bar m) - \tfrac{i}{\hbar}\int \varphi\,\Delta m\Big\}, \qquad \varphi = 4\pi G\, G_{\mathrm{ret}} * A. \tag{4}$$

For a fixed record, W is a positive-definite kernel on path pairs on a region if and only if

$$K \equiv D_0 - \tfrac{1}{4} D_1^{\mathsf{T}} D_2^{-1} D_1 \succeq 0 \tag{5}$$

on that region: the decoherence–diffusion trade-off [4, 11]. At $\omega=0$ it states that a classical mediator's noise is at least its full response, $D \geq |\chi|/\hbar$, the static form of Tilloy and Diósi [6].

2.3 From regions to momenta

The trade-off (5) is a statement about regions. The floors below need it momentum by momentum.

Lemma 2.1 (Per-momentum trade-off)**.** *Let D_0, D_1, D_2 be translation invariant on Minkowski space, with locally integrable, polynomially bounded symbols on the physical record space, and let the dynamics be CP on every past region $R = \{t_i < t < T\}$, with kernels on R the restrictions of the global ones. Then, for almost every p, near-null p included,*

$$\tilde{D}_0(p) \;\succeq\; \tfrac{1}{4}\,\tilde{D}_1(p)^\dagger\,\tilde{D}_2(p)^{-1}\,\tilde{D}_1(p). \tag{6}$$

Proof. For $A \succ 0$, $\langle \chi, A^{-1}\chi\rangle = \sup_\psi [2\langle\chi,\psi\rangle - \langle\psi, A\psi\rangle]$. CP on R therefore gives $\langle f, D_0 f\rangle \geq \frac{1}{4}(2\langle D_1 f,\psi\rangle - \langle\psi, D_2\psi\rangle)$ for every $\psi \in C_c^\infty(R)$ and every f supported in R; only $D_1 f$ on $\operatorname{supp}\psi$ enters. For fixed f, translation invariance lets R be shifted and enlarged to contain any compact set, so the inequality holds for every $\psi \in C_c^\infty(\mathbb{R}^4)$. C_c^∞ is dense in the form domain of a tempered, locally integrable symbol and Plancherel holds in any signature, so the supremum is $\frac{1}{4}\int |\tilde{D}_1\tilde{f}|^2/\tilde{D}_2\, d^4p/(2\pi)^4$. Wave packets concentrated at p give (6). □

Remark 2.1. On any *fixed* region the compressed kernel satisfies $(D_2|_R)^{-1} \prec (D_2^{-1})|_R$, and near the cone the gap is large; CP on one region does not imply (6). CP on every region does. A packet at $u = 1\ \mathrm{m}^{-2}$ on the 140 keV shell needs a null extent $k/u = 7.1 \times 10^{11}$ m, far below the Hubble length, so the finite age of the universe does not reopen the gap. A kernel that is CP region by region but not translation invariant carries an origin or a frame; it belongs to the shielded class of Section 6.

The same steps at one momentum give the response form of the trade-off. With $S(J) = \langle J, D_0 J\rangle + c^2\langle GJ, D_2 GJ\rangle$ the total noise on a source J, $c = 4\pi G/\hbar$ and G the propagator of the mean field, Cauchy–Schwarz in the D_2 inner product and the arithmetic–geometric mean inequality give

$$S(J) \;\geq\; c\,|\langle D_1 J, GJ\rangle|, \tag{7}$$

whatever the form of D_2. In units (3) this reads $s \geq |f|/u$ for a scalar source with a retarded drift (retarded and advanced agree at spacelike p). Static Newton fixes $|f(u)| = 1$ on $u \in [1, 10^6]\ \mathrm{m}^{-2}$.

3 The scalar strip

3.1 The covariant diffusion and the conformal mode

Linearised about Minkowski space, covariant CQ gravity gives the metric perturbation h the weight [11, 12, 13]

$$\exp\Big\{-\tfrac{1}{2}\int (\alpha R_{\mu\nu}R^{\mu\nu} - \beta R^2)\Big\}, \qquad b \equiv \beta/\alpha < \tfrac{1}{3}, \tag{8}$$

a diffusion with DeWitt form $D_2^{-1} = \alpha I - \beta\eta\eta$. The explicit decoherence term is $\frac{1}{8}\Delta T\, D_0\, \Delta T$ with $D_0 \succeq 64\pi^2 G^2 D_2^{-1}$, saturated by the authors. Matter couples through $\frac{1}{2}h_{\mu\nu}T^{\mu\nu}$. In the scalar–vector–tensor split ψ and h^{TT} are dynamical and the Bardeen potential Φ and the vector

V_i are constraints [13]. The scalar correlator is [13]

$$\langle \Phi\Phi \rangle = \frac{3}{2\alpha k^4} + \frac{1}{\alpha p^4}\Big[\frac{\omega^2}{k^2} + \frac{-3\alpha + 10\beta}{4(\alpha - 3\beta)}\Big]. \tag{9}$$

We rederived (9) from the quadratic form of (8), with the four gauge modes removed and the Bardeen functionals $\Phi = \frac{1}{2}h_{00} + \frac{\omega}{k}\hat{k}\cdot h_0 + \frac{3\omega^2}{4k^2}(\hat{k}\hat{k} - \frac{1}{3}\delta)\!:\!h$, $\psi = \frac{1}{4}(\delta - \hat{k}\hat{k})\!:\!h$.

Proposition 3.1 (Near-cone conformal mode)**.** *On the strip $u \ll k^2$,*

$$\langle \Phi\Phi \rangle = \frac{c_d(b)}{\alpha u^2} - \frac{1}{\alpha k^2 u} + O(1), \qquad c_d(b) = \frac{1-2b}{4(1-3b)}, \tag{10}$$

with $c_d \in [\frac{1}{6}, \frac{1}{2}]$ for $b \le \frac{1}{4}$. At every p, $\langle(\Phi+\psi)^2\rangle = 3/(2\alpha k^4)$, so that

$$\frac{\langle(\Phi+\psi)^2\rangle}{\langle \Phi^2\rangle} = \frac{3}{2c_d}\Big(\frac{u}{k^2}\Big)^2[1 + O(u/k^2)] \le 9\Big(\frac{u}{k^2}\Big)^2[1 + O(u/k^2)]. \tag{11}$$

Near the cone $h_{\mu\nu} \approx 2\psi\eta_{\mu\nu}$: a conformal mode.

Proof. At $p^4 = u^2$ and $\omega^2/k^2 = 1 - u/k^2$ the bracket of (9) is $1 + \frac{-3+10b}{4(1-3b)} - \frac{u}{k^2} = \frac{1-2b}{4(1-3b)} - \frac{u}{k^2}$, which gives (10). For $\Phi + \psi$: on $h = 2\sigma\eta$ the functionals give $\Phi = -\sigma$, $\psi = \sigma$, so $\Phi + \psi$ is gauge and Weyl invariant. At quadratic order $R_{\mu\nu}R^{\mu\nu} = \frac{1}{2}C_{\mu\nu\rho\sigma}C^{\mu\nu\rho\sigma} + \frac{1}{3}R^2$ up to a total derivative, so the weight is $\frac{\alpha}{2}C^2 + \alpha(\frac{1}{3} - b)R^2$ with both coefficients positive for $b < \frac{1}{3}$. On the scalar sector, where $\Phi + \psi$ lives and the weight is positive, the variance of a linear functional L is $\sup_h (Lh)^2/(\text{weight})$; over all h the supremum would be infinite, since the vector weight is negative at spacelike p (Section 6.2). Since $R(2\sigma\eta) \propto p^2\sigma \neq 0$ off the cone, every h can be shifted along the conformal direction to $R = 0$ without changing $L = \Phi + \psi$ or C; the supremum is therefore attained at $R = 0$, and it does not depend on b. Evaluating the quadratic form gives $3/(2\alpha k^4)$ (checked numerically to 2×10^{-9} over $b \in [-100, 0.33]$ and $\omega/k \in [0, 3]$). Dividing by (10) and using $c_d \ge \frac{1}{6}$ for every $b < \frac{1}{3}$ gives (11). □

3.2 Coupling: rest mass, not Maxwell

On a conserved near-null source (Section 3.4)

$$\tfrac{1}{2}h_{\mu\nu}J^{\mu\nu} = (\Phi + \psi)\rho - \psi\rho_m \;\approx\; \Phi\,\rho_m, \tag{12}$$

so the mode couples to the rest-mass density $\rho_m = -T^{\mu}{}_{\mu}$; the traceless electromagnetic binding fields drop out. A slow particle has geodesic acceleration $\mathbf{a} = -\nabla\Phi$. Maxwell's equations are conformally invariant in four dimensions, so the photon field does not see the mode: in conformal variables Maxwell is flat and the rest mass varies, a slow charge has coordinate acceleration $-\nabla\Phi$ and radiates by Larmor, with no Coulomb-scattering term to cancel it. The non-conformal remainder is $O(u/k^2)$, 3.9×10^{-16} at 10 keV at the top of the band, $u = 10^6\ \mathrm{m}^{-2}$. Near the cone the scalar sector of covariant CQ gravity acts on slow matter as a stochastic Nordström field.

3.3 The mean field is shielded; the noise is not

General relativity's retarded response of Φ to any conserved source (ρ, τ_{ij}) is

$$\frac{\Phi_{\text{resp}}}{16\pi G} = \frac{1}{k^2}\Big[-\frac{\rho - \operatorname{tr}\tau}{4} + \frac{\rho u}{2k^2}\Big], \tag{13}$$

with no pole at $u = 0$: the instantaneous Poisson constraint at every ω. The $1/u$ response sits in the tensor channel (Section 5). The response form (7) therefore gives slow matter no near-cone floor. The floor comes from the diffusion structure (8), through both channels at once.

3.4 The conserved near-null source

At near-null p every conserved source has the form

$$J = A\, ll + B\,(le + el) + C_{ab}\, e_a e_b, \qquad l = (1, \tfrac{\omega}{k}\hat{k}), \quad l \cdot l = -\frac{u}{k^2}, \quad e_a \perp \hat{k}, \tag{14}$$

and conservation forces $T^{0z} \approx T^{00}$. For a slow atom the relevant source is the transition element $\langle f\gamma|T^{\mu\nu}(p)|0\rangle$ of the full emission process. It includes the photon and the Maxwell stress of the atom's Coulomb field and is conserved because the total T is, so (14) is a property of the amplitude, not of the kinetic stress of the atom. The trace sits in the transverse block, $\operatorname{tr} C = T + Au/k^2 \approx -\rho_m$. Write $\kappa \equiv C{:}C/\rho_m^2 \geq \frac{1}{2}$, with equality when the traceless part C^{TT} vanishes. The saturated explicit kernel per α splits as

$$\frac{N}{\alpha} = J{:}J - bT^2 = (\tfrac{1}{2} - b)(\operatorname{tr} C)^2 + C^{\text{TT}}{:}C^{\text{TT}} - 2B^2\frac{u}{k^2} + (1-b)A^2\frac{u^2}{k^4} + 2bA\operatorname{tr} C\,\frac{u}{k^2}. \tag{15}$$

The ll part, which carries T^{00}, T^{0z} and T^{zz} together, has weight $(1-b)u^2/k^4 \leq (1-b)\,1.5\times10^{-27}$ for $u \leq 10^6$ m^{-2} and $k \geq k(1\text{ keV}) = 5.07\times10^9$ m^{-1}: the energy density carries none of the floor. The trace enters with weight $\frac{1}{2} - b > 0$, the anisotropic stress with weight one, and the B term, the non-positive direction of the DeWitt form, at $O(u/k^2)$.

3.5 The floor

Theorem 3.2 (Scalar floor)**.** *Let the CQ dynamics be Gaussian, translation invariant and CP on every past region, with matter coupled through $\frac{1}{2}h_{\mu\nu}T^{\mu\nu}$ and Newton's force carried by the classical field. Let the diffusion weight on physical modes be $D_2^{-1} = a_2(p^2)P^{(2)} + a_0(p^2)P^{(0s)}$ with $a_2, a_0 > 0$, the DeWitt form (8) included. Then on the conserved near-null source (14)*

$$s(u) \;\geq\; \frac{2}{3}\,\frac{|f(u)|}{u}\,[1 + O(u/k^2)] \tag{16}$$

for every α, every $b < \frac{1}{3}$ and every p^2-dependent regulator or running of a_2, a_0.

Proof. For the DeWitt weight the two channels on the same source are

$$s_{\text{exp}} = 4\pi\ell_{\text{P}}^2\,\alpha(\kappa - b)\,|f|^2, \qquad s_{\text{dyn}} = \frac{\kappa - \frac{1-4b}{4(1-3b)}}{4\pi\ell_{\text{P}}^2\,\alpha\,u^2}, \tag{17}$$

the first from (15) with D_0 saturating (6) (Lemma 2.1), the second from (10) on the source.

Both increase with κ, so their product is least at $\kappa = \frac{1}{2}$:

$$s_{\exp}\, s_{\mathrm{dyn}} \geq (\tfrac{1}{2} - b) c_d \frac{|f|^2}{u^2} = \frac{\tilde{g}^2 |f|^2}{4u^2}, \qquad \tilde{g}^2 = \frac{(1-2b)^2}{2(1-3b)} \geq \frac{4}{9}, \tag{18}$$

with the minimum at $b = \frac{1}{6}$. The product does not contain α, and $s \geq 2\sqrt{s_{\exp} s_{\mathrm{dyn}}} = \tilde{g}|f|/u \geq \frac{2}{3}|f|/u$. For general positive (a_2, a_0) the invariant form on conserved sources is $\frac{1}{4} J\langle hh\rangle J = (x_2/a_2 + x_0/4a_0)/p^4$ with $x_2 = J \cdot P^{(2)} \cdot J$ and $x_0 = T^2/3$, and Cauchy–Schwarz gives

$$(a_2 x_2 + a_0 x_0)\Big(\frac{x_2}{a_2} + \frac{x_0}{4a_0}\Big) \geq \big(x_2 + \tfrac{1}{2} x_0\big)^2 \geq \frac{\rho_m^4}{9}, \tag{19}$$

since at $\kappa \geq \frac{1}{2}$, $x_0 = \rho_m^2/3$ and $x_2 \geq \rho_m^2/6$. The DeWitt weight at $b = \frac{1}{6}$ attains the bound. □

Remark 3.1. (i) The floor (16) is not the square of a response: the scalar response is shielded (13). It needs a weight diagonal in the Lorentz-invariant projectors. A frame in the scalar weight breaks this structure; Section 6.5 lists what that costs. (ii) At $b \to \frac{1}{3}^-$ the factor $\tilde{g}^2$ rises monotonically from its minimum and diverges; at $b = \frac{1}{3}$ the scalar is Weyl-degenerate and unweighted, and $b > \frac{1}{3}$ is not positive. Nothing opens at the edge. Hirotani and Matsumura report that a strain-noise term vanishes at $\beta = \frac{1}{3}$ [19]; that factor comes from a non-conserved stochastic source in harmonic gauge, and on the conserved source (14) the physical scalar weight is $a_0 = \alpha - 3\beta$. (iii) At $\omega = 0$ conservation gives $\kappa = 1$ and the floor is $\sqrt{(1-b)(3-8b)/(1-3b)}/k^2 \geq \frac{5}{3}/k^2$, above the static mediator value $1/k^2$. This static floor is Oppenheim and Sajjad's [13]. With $x = 1 - 3b$, $(1-b)(3-8b)/(1-3b) = (1+8x)(2+x)/(9x)$ is their ratio of the scalar channel's product of diffusion and decoherence to the tensor channel's, whose minimum 25/9 at $b = \frac{1}{6}$ they give. The noise floor is its square root, because the least noise is twice the geometric mean of the two channels.

3.6 The integrated floor

Newton's law is tested with $|f| = 1$ from 1 mm to 1 m, that is $u \in [1, 10^6]$ m^{-2}. By Lorentz invariance of the scalar weights this band is a near-cone strip at every photon energy, and

$$\int s\, du \;\geq\; \frac{2}{3} \ln 10^6 = 9.21. \tag{20}$$

For a theory whose classical field couples to the trace alone (Nordström-type CQ gravity [16]) the response form (7) applies directly, $s \geq |f|/u$, and $\int s\, du \geq \ln 10^6 = 13.8$.

4 X-ray conversion and the XENONnT exclusion

4.1 The emission rate

For Gaussian noise on the mass density with $\langle\varphi\varphi\rangle(\mathbf{q}) = \hbar^2 D(\mathbf{q})$, every particle receives the mass-independent acceleration $-\nabla\varphi$. Closure over atomic final states gives the photon spectrum per atom

$$\frac{d\Gamma}{dE} = \frac{e^2\hbar}{4\pi^2 \varepsilon_0 c^3 E} \int \frac{d^3q}{(2\pi)^3}\, D(\mathbf{q})\, q^2 \left\langle \sin^2\!\angle(\mathbf{q}, \mathbf{k})\; S(|\mathbf{q} - \mathbf{k}|)\right\rangle_{\hat{k}}, \qquad k = \frac{E}{\hbar c}, \tag{21}$$

with $S(Q) = |Z - F(Q)|^2 + \sum_a n_a(1 - f_a^2)$ the charge structure factor of the atom. This is Manti *et al.*'s rate for a stationary correlation [21] evaluated on the near-cone correlation; for smooth noise it reduces to the Diósi–Penrose pair formula used in the published bounds [8, 10]. There is no atomic cancellation at keV: a shaken neutral atom radiates through $Z - F(Q)$, which vanishes only as $Q \to 0$.

Near-cone noise has $|\mathbf{q}| = k$ and is isotropic, so the structure factor is sampled at the Thomson momentum transfer $Q = 2k\sin(\theta/2)$. Relative to the uncorrelated value $Z^2 + Z$,

$$A_{\mathrm{cone}}(E) = \frac{\langle \frac{3}{2}\sin^2\theta\; S(2k\sin\frac{\theta}{2})\rangle}{Z^2 + Z}. \tag{22}$$

For an isolated local-density xenon atom:

E (keV)	1	10	30	50	140
A_{cone}	0.002	0.258	0.633	0.786	0.955

The emission integral $I(E) = 2\pi\int(u + E^2)^{3/2}s\,du \geq 2\pi E^3\int s\,du$ makes the E^2A_{cone} template a lower bound on the signal.

4.2 The trace converts like the conformal mode

The floor of Theorem 3.2 sits on the explicit channel's coupling to $\mathrm{tr}\,C = -\rho_m$. Noise on that trace is the transverse-trace metric $h_{ab} = \chi\delta_{ab}$ $(a, b \perp \hat{q})$, whose Bardeen potentials are

$$\Phi = -\frac{\chi\omega^2}{2q^2}, \qquad \psi = \frac{\chi}{2}, \qquad \Phi + \psi = \frac{\chi u}{2q^2}, \tag{23}$$

the conformal mode near the cone. For a metric mode acting on charges at rest with charge form factor $R(Q)$, linearised geodesics and the linearised constitutive relations of the Coulomb field give a radiated field $\propto \hat{k} \times \mathbf{J}$ with

$$\mathbf{J} = R(Q)\Big[\mathbf{v} + \omega\,\frac{\Delta\varepsilon\cdot\mathbf{Q}}{Q^2} - \frac{\mathbf{k}\times(\mathbf{g}\times\mathbf{Q})}{Q^2}\Big], \qquad \Delta\varepsilon = (\tfrac{1}{2}h_{00} + \tfrac{1}{2}\,\mathrm{tr}\,h)\mathbb{1} - h_{ij}, \quad g_i = h_{0i}, \tag{24}$$

$\mathbf{v}$ the geodesic velocity. The last two terms are the Maxwell stress of the atom's own Coulomb field. $\mathbf{J}$ is invariant under $h \to h + i(p\zeta + \zeta p)$ for any ζ and off-shell $\mathbf{q}$. For the trace mode (23) it equals the Larmor amplitude of the conformal kick up to a relative remainder $\leq 3u/Q^2$. In the gauge where the noise couples only to stresses, the photon comes from the Coulomb field's Maxwell stress, which is of order one, not from the particles' kinetic stress, which is of order v^2/c^2. Because the Maxwell stress tensor is traceless, the trace the floor uses is the rest-mass trace whether or not the noise couples to the Coulomb field. The floor and the field noise share one conversion, A_{cone}.

4.3 The limit

We use a profile likelihood of the XENONnT SR0 spectrum [10]: 84 bins from 1 to 140 keV, 8 profiled background components, an efficiency nuisance, 1.16 t y, one-sided 95% CL. Without calibration it reproduces the published Diósi–Penrose limit at $C = 1$ as $R_0 > 4.483$ Å against the published 4.5 Å. The margins below use the uncalibrated likelihood; the calibration factor

Table 1. *XENONnT margins: lower bound on $\int s\,du$ divided by the limit (25). Rows 1 and 4 need neither α nor any coherence bound; row 2 needs neither the trade-off nor the explicit-channel vertex.*

	Row, applies to	$\int s\,du \geq$	Margin
1	Floor (16) on $u \in [1, 10^6]$ m^{-2}: every α, $b < \frac{1}{3}$, every invariant (a_2, a_0)	9.21	1.19×10^8
2	Field noise alone, constant couplings at the coherence edge $D_{2,\mathrm{scalar}} \geq 10^{-71}$	1.9×10^{-4}	2.5×10^3
3	Invariant ε prescription, $\langle\Phi\Phi\rangle = D_G/(p^4+\varepsilon^2)$, $D_G \geq 10^{-71}$	3.05×10^{-3}	3.9×10^4
4	Trace-coupled (Nordström-type) CQ, $s \geq \lvert f\rvert/u$	13.8	1.79×10^8

1.0115 that closes the gap would raise every margin by that factor. For the $E^2 A_{\mathrm{cone}}$ template, with A_{cone} interpolated through the five tabulated points,

$$C \int s\,du \;<\; 7.7 \times 10^{-8} \qquad \text{(without the atomic factor: } 7.4 \times 10^{-8}\text{)}. \tag{25}$$

4.4 The margins

Table 1 collects the margins. Row 1 is (20). Row 2 integrates the field noise $c_d/(\alpha u^2)$ from $u = 1$ m^{-2}: $\int s\,du = (c_d/\alpha)/(4\pi\ell_{\mathrm{P}}^2 \cdot 1\ \mathrm{m}^{-2})$. At fixed scalar diffusion constant $D_{2,\mathrm{scalar}} = (3\alpha - 8\beta)/(\alpha(\alpha - 3\beta))$, bounded below by matter-wave coherence [14, 13], $\min_b(c_d/\alpha)/D_{2,\mathrm{scalar}} = \frac{1}{16}$, reached as $b \to -\infty$; at $b = 0$ the margin is larger. Row 3 is the regulated form of [12] at the same coherence edge. Row 4 is the response form (7) for a scalar source.

Varying the inputs: the elastic structure factor alone gives 1.18×10^8 for row 1. An envelope for the order-v^2/c^2 correction of bound electrons to the coherent amplitude, $c_v \in [-3, 3]$ times the kinetic form factor, spans 1.11–1.27 $\times 10^8$. We quote row 1 as $\geq 1.1 \times 10^8$.

4.5 The boost the exclusion needs

The full margins use the strip up to 140 keV, which maps $u = 1$ m^{-2} onto the photon shell with a boost $k/\sqrt{u} = 7.1 \times 10^{11}$. Exclusion needs far less. If the scalar sector is Lorentz invariant only up to $\gamma_{\max}$, the usable strip at energy E is $u \geq (k/\gamma_{\max})^2$, and the refit gives:

$\gamma_{\max}$	10^9	10^8	1.2×10^7	10^7
floor margin	8.3×10^6	1.6×10^4	2.1	0.68

The floor margin reaches one at $\gamma_{\max} = 1.1 \times 10^7$; the constant-coupling window (row 2) at $\gamma_{\max} = 1.6 \times 10^{10}$.

5 The tensor strip: a second exclusion

The scalar floor needs a diffusion weight that is diagonal in the Lorentz-invariant projectors, and it needs Lorentz invariance to carry Newton's tested band to the light cone. A frame in the scalar weight removes it (Section 6.5). This section shows that the tensor sector gives an exclusion that uses neither step.

5.1 The response is a tensor

Normalised so that a static density gives Newton's response, $\chi = \rho^2/k^2$, the response of general relativity to a conserved source is

$$\chi(J) = \frac{2\,(J{:}J - \frac{1}{2}T^2)}{u}. \tag{26}$$

On the near-null source (14), the split (15) at $b = \frac{1}{2}$ gives

$$J{:}J - \tfrac{1}{2}T^2 = C^{\mathrm{TT}}{:}C^{\mathrm{TT}} - 2B^2\frac{u}{k^2} + \frac{1}{2}A^2\frac{u^2}{k^4} + A\,\mathrm{tr}\,C\,\frac{u}{k^2}, \qquad \chi(J) = \frac{2\,C^{\mathrm{TT}}{:}C^{\mathrm{TT}}}{u} + O(k^{-2}). \tag{27}$$

The light-cone pole that the scalar response lacks (13) is entirely in the transverse-traceless block. It is the propagator of the theory's own mean field: in the covariant theory the classical Einstein equation; in a theory with a preferred foliation, $1/(c_T^2k^2 - \omega^2)$ with $|c_T - 1| \lesssim 10^{-15}$ [26]. For $c_T \neq 1$ the strip sits on the spacelike side of the tensor cone, where the response is real, and the near-null relations (14) hold to $O(|c_T - 1|)$.

5.2 The floor holds for every weight

Theorem 5.1 (Tensor floor)**.** *Assume*

(T1) Gaussian CQ dynamics, translation invariant and CP on every past region;

(T2) a drift whose response on conserved sources at the near-null momentum p is the tensor response (27);

(T3) matter coupled through $\frac{1}{2}h_{\mu\nu}T^{\mu\nu}$ to the total conserved $T^{\mu\nu}$, the atom's Coulomb field included.

Then for every positive diffusion kernel D_2, Lorentz invariant or carrying any frame, the matter noise on every unit transverse-traceless polarisation e obeys, in units (3),

$$s_e(u) \;\geq\; \frac{2}{u}. \tag{28}$$

Proof. Lemma 2.1 uses only (T1) and gives (6) at p. The response form (7) follows from it by Cauchy–Schwarz in the D_2 inner product and needs only $D_2 \succ 0$. For $J = e$, $e{:}e = 1$, it gives $s_e \geq |\chi(e)| = 2/u$ by (27). □

Remark 5.1. Theorem 3.2 and Theorem 5.1 differ in what they square. The scalar floor is not the square of a response, because the scalar response is shielded; it is the product of two channels whose weights must be diagonal in the invariant projectors. The tensor floor is the square of a response, and Cauchy–Schwarz holds in any positive inner product. A frame in the weight changes which inner product is used, not the bound.

5.3 Conversion of tensor noise

The tensor noise is a random metric field on the transverse block, $h_{ab} = \sum_e \xi_e e_{ab}$, with orthonormal polarisations e and variances N_e. Through the amplitude (24), a tensor mode acts on an atom at rest through the Maxwell stress of its Coulomb field; the particles' kinetic stress

gives only order v^2/c^2. Per unit variance on orthonormal components the angular weights of $|\hat{k}\times\mathbf{J}|^2$ are $\sin^2\theta/8$ for the trace and, for the two tensor polarisations summed and averaged over azimuth, $\sin^2\theta\,(1-\frac{1}{2}\sin^2\theta)/(4(1-\cos\theta)^2)$; after the azimuthal average there are no trace–tensor cross terms. With the elastic structure factor the ratio of tensor to trace emission is

$$X_{\mathrm{TT}}(E)=\frac{2\big\langle \sin^2\theta\,(1-\frac{1}{2}\sin^2\theta)\,S_{\mathrm{el}}(Q)/(1-\cos\theta)^2\big\rangle}{\big\langle \sin^2\theta\,S_{\mathrm{el}}(Q)\big\rangle},\qquad Q=2k\sin\tfrac{\theta}{2},\tag{29}$$

E (keV)	5	10	30	140
X_{TT}	1.62	2.29	4.27	9.60

$S_{\mathrm{el}}=|Z-F|^2$ is the exact ground-to-ground factor of the atom, so the elastic conversion needs no closure; the inelastic channels only add to it, so the elastic value is the lower side. The conversion is forward peaked. In liquid xenon inter-atomic interference acts at $Q<2$ Å^{-1}, which holds 23%, 7%, 1.6% and 0.5% of the numerator of (29) at 5, 10, 30 and 140 keV; it is immaterial at the energies that set the limit.

Proposition 5.2 (Tensor emission floor)**.** *Under Theorem 5.1 the photon rate from tensor noise is at least the rate of trace noise with*

$$s_{\mathrm{eq}}(u)=\frac{X_{\mathrm{TT}}(E)}{u}=\frac{3}{2}\,X_{\mathrm{TT}}(E)\times\frac{2/3}{u}.\tag{30}$$

Proof. Let $N_{ee'}$ be the noise covariance on orthonormal polarisations, of any form. The coupling $\frac{1}{2}h{:}J$ gives the decoherence $\lambda\sum_{ee'}N_{ee'}J_eJ_{e'}$ with $J_e=e{:}J$ and $\lambda=\frac{1}{4}$. The trace polarisation is $e_{\mathrm{tr}}=\delta_{ab}/\sqrt{2}$, so $J_{\mathrm{tr}}=\mathrm{tr}\,C/\sqrt{2}=-\rho_m/\sqrt{2}$ and $s=\lambda N_{\mathrm{tr}}/2$: trace noise of strength s has variance $2s/\lambda$. In the emission, the azimuthal integral over photon directions removes every cross conversion identically, whatever the azimuthal structure of N: the trace–tensor term is $\propto\cos\theta\sin\theta\cos2\phi$ and the $+\times$ term $\propto\sin^2\theta\cos2\phi\sin2\phi$, as the helicities 0 and ±2 about $\hat{q}$ require. Only the diagonal N_{ee} enter, and by (28) $\lambda N_{ee}\geq 2/u$ for each tensor polarisation. The tensor emission is therefore at least $(2/\lambda u)$ times the summed tensor conversion, which is X_{TT} times the trace conversion per unit variance. That equals the emission of trace noise with $s=X_{\mathrm{TT}}/u$. □

5.4 The margin

The template is the one of Section 4 with A_{cone} restricted to its elastic part and multiplied by $X_{\mathrm{TT}}(E)$, refitted to the same likelihood. The product $A^{\mathrm{el}}_{\mathrm{cone}}X_{\mathrm{TT}}=3\langle\sin^2\theta(1-\frac{1}{2}\sin^2\theta)S_{\mathrm{el}}/(1-\cos\theta)^2\rangle/(Z^2+Z)$ cancels the trace conversion, so the tensor template uses no conformal-mode input. Two bands:

1. *Newton's band*, $u\in[1,10^6]$ m^{-2}, the band of the scalar floor. The margin is 1.7×10^9 (median expected 9.8×10^8). Divided by the elastic scalar margin 1.18×10^8 and by $\frac{3}{2}$, it gives a template-weighted $X_{\mathrm{TT}}=9.4$; that number is an output of the refit, not an input.

2. *The tensor's own band.* The tensor response is the theory's propagator, not a static test carried to the cone, so the band need not stop at 10^6 m^{-2}. On $u\in[1\text{ m}^{-2},10^{-2}k^2]$, where the near-null relations hold to 1%, the margin is 6.0×10^9; at 140 keV the logarithm is 50.0 instead of 13.8. We keep the lower end at 1 m^{-2}. Below it the wave packets of

Lemma 2.1 approach the Hubble length; at the cosmological friction scale $u_\eta = 2kH_0/c = 1.0 \times 10^{-14}$ m^{-2} (140 keV) they reach it.

Both bands use the spacelike side of the pole only. On the timelike side $|\chi| = 2/|u|$ as well, so including it would double every tensor margin; we leave that factor out to keep the same one-sided band as the scalar strip. Setting the normalisation $\frac{3}{2}$ to one and freezing $X_{\rm TT}$ at its 5 keV value gives 1.9×10^8. That number tests sensitivity to those two inputs only; it is not a floor against the real uncertainties, which are the atom model and (T3).

Remark 5.2 (What is an estimate). Theorem 5.1 is exact under (T1)–(T3), and the normalisation of Proposition 5.2 follows from the coupling. The margin is an estimate for two reasons. (i) $X_{\rm TT}$ uses the elastic structure factor of an isolated local-density atom and the near-cone kinematics. The elastic factor is exact for that atom and needs no closure, and the inelastic part only adds; a full calculation would raise the margin, not lower it. (ii) Assumption (T3) is about how classical gravity couples to field energy. The equivalence principle fixes the coupling to the Coulomb field's energy, which is part of every atom's mass. Spatial diffeomorphisms, which a preferred foliation keeps, give only $\partial_0 T^{0j} + \partial_i T^{ij} = 0$; for a static atom $T^{0j} = 0$, so the total stress it couples must be divergence-free. Kinetic and Coulomb stress separately have divergence $\mp$ the Lorentz force density, so keeping one without the other breaks the symmetry a foliation keeps. A model could still drop the Coulomb stress. The tensor mode then reaches the atom only through the particles' kinetic stress. In the gauge where the noise is h_{ab} alone, it changes an electron's velocity by $\delta\mathbf{v} = -h \cdot \mathbf{v}$, an incoherent current whose elastic part vanishes. With closure over final states and isotropic orbitals, the two polarisations give $|\hat{k} \times \mathbf{J}|^2 = \frac{2}{3}\sum\langle v^2/c^2\rangle(1 - \frac{1}{2}\sin^2\theta)$ per unit variance, with $\sum\langle v^2/c^2\rangle = 0.77$ for the local-density xenon atom. The tensor template becomes the constant $\frac{16}{3}\sum\langle v^2/c^2\rangle/(Z^2 + Z) = 1.4 \times 10^{-3}$ in place of $A^{\rm el}_{\rm cone} X_{\rm TT}$, and the Newton-band margin is 2.6×10^5. The estimate keeps the floor $\lambda N_{ee} \geq 2/u$, which was derived for conserved sources. This is a rough estimate: without (T3) the source is not conserved and the result depends on the gauge in which the coupling is written. On this estimate, dropping the Coulomb field alone does not escape; a coupling blind to the atom's stress altogether would.

5.5 What the two exclusions share

The two exclusions share Lemma 2.1, the emission rate (21), the local-density xenon atom, the coupling $\frac{1}{2}h_{\mu\nu}T^{\mu\nu}$, linearisation, the Gaussian class, translation invariance, and the XENONnT likelihood. Each also assumes that the mean field of its own sector is that of general relativity. The tensor exclusion does not use the projector-diagonal weight of Theorem 3.2, the Lorentz transfer of Newton's band (no boost γ appears), or the conformal-mode conversion. In exchange it assumes that general relativity's tensor propagator holds at keV momenta. Gravitational-wave observations test that propagator only at photon-equivalent energies more than fifteen orders of magnitude lower, so this is an assumption, not a measured fact. The exclusions are independent in this sense: different escapes defeat each.

Escape	Scalar exclusion	Tensor exclusion
A frame in the scalar weight	removed	survives
A non-perturbative scalar mass	removed	survives
A tensor sector off the light cone at keV	survives	removed
A coupling blind to the atom's stress	survives	removed
Lemma 2.1, the atom model or the Gaussian class	removed	removed

5.6 What it excludes, and what escapes

Under (T1)–(T3) and a tensor sector that is that of general relativity at keV momenta, the tensor floor excludes the published covariant theory, both completions proposed by its authors (a unit timelike vector in the scalar and vector weights, and in the vector weight only [13]), and classical CQ gravity with a dynamical preferred foliation. In the last case the foliation removes the scalar strip, because the Bardeen potential's response in the preferred frame has no light-cone pole, but the tensor propagator keeps its pole with $c_T = 1$ to 10^{-15}.

The tensor floor is escaped only by leaving (T2) or (T3):

- a tensor dispersion that leaves the light cone below the X-ray scale. Take a Lifshitz-type dispersion $\omega_T^2(q) = q^2 + q^4/k_*^2 + q^6/k_*^4$ with $M_T = \hbar c k_*$. At photon frequency ω the pole moves to $q_T < k$, and the response becomes $2/|D|$ with $D = \omega_T^2(q) - \omega^2$. Two factors change. The measure is $4\pi q^2 dq = 4\pi q_T^2\, dD/(d\omega_T^2/dq)$. The tensor amplitude at q_T carries no force factor: a transverse-traceless mode has $h_{00} = h_{0i} = 0$, so $\mathbf{v} = 0$ in (24), and it radiates through the Coulomb stress, $\mathbf{J} = R\,\omega^2 (e \cdot \hat{k})/Q^2$ with $|\mathbf{Q}| = |\mathbf{q} - \mathbf{k}| \approx k$ for $q_T \ll k$. Relative to the light cone the template is multiplied by

$$\frac{q_T^2\, T(q_T)/(d\omega_T^2/dq)}{(k/2)\, T(k)}, \qquad T(q) = \Big\langle \sin^2\theta\,(1 - \tfrac{1}{2}\sin^2\theta)\, S_{\rm el}(Q)\, \frac{k^4}{Q^4} \Big\rangle, \tag{31}$$

  with $Q^2 = k^2 + q^2 - 2kq\cos\theta$. It tends to $k_*^2 T(q_T)/(3k^2 T(k))$ deep in the Lifshitz regime. At $M_T = 10$ eV it is 4.4×10^{-7}, 3.4×10^{-8} and 8.2×10^{-10} at 10, 30 and 140 keV. Refitting the moved-pole template to the same likelihood, the Newton-band margin is 12 at $M_T = 30$ eV and 136 at 100 eV, and it falls to one at $M_T = 8.6$ eV (median expected 9.1 eV). A classical theory of this kind needs $M_T \lesssim 10$ eV. That edge is for the $z = 3$ dispersion above with unit coefficients. With the q^4 term alone ($z = 2$) the deep suppression is $(k_*/k)^{3/2}/2$ instead of $k_*^2/(3k^2)$, and the edge falls to $M_T = 0.24$ eV. The model-independent condition is $q_T^2/(d\omega_T^2/dq) \lesssim 10^{-9}\, k/2$ over 50–140 keV. In the healthy extension of Hořava gravity the helicity-2 dispersion is set by operators independent of those that fix the lapse [24, 25], so M_T is a separate scale from the one that regulates Newton's force;

- a form factor in the tensor drift at keV momenta;

- a matter coupling blind to the stress of the atom, Coulomb field and particles alike (Remark 5.2);

- quantized gravitons with only the Newtonian constraint classical, which is a hybrid, not classical gravity;

- jump-class (non-Gaussian) CQ dynamics [15, 27] whose keV emission arrives in multi-photon bursts (Section 6);

- a frame carried by the matter coupling itself, a form factor in the matter four-velocity that acts on the inelastic vertices where X-ray emission lives (Section 6.5).

A scalar mass, which Section 6 leaves open for the scalar strip, does not act on the tensor sector and does not escape the tensor floor.

6 Objections and scope

6.1 X-ray bounds constrain white noise only

Curceanu *et al.* stress that X-ray bounds constrain Markovian, effectively white noise, and that suppressing the spectrum at high frequency can weaken them [22]; Vinante and Ulbricht make the same point for heating [23]. The caveat concerns spectra suppressed at high frequency. The near-cone spectrum is fixed by the theory: at frequency ω the force-noise density is $\int d^3q\, q^2 S \propto \omega^3 \int s\, du$, which is blue, and Lorentz invariance places its support at keV. The emission rate (21) uses the spectral density at $\omega = E/\hbar$, which is what Section 4 computes. In Oppenheim and Sajjad's analysis the matter decoherence obtained after integrating out ψ is non-Markovian [13]. That does not enter: $s(u)$ is the field noise of ψ with its own propagator, and the golden-rule rate needs only its spectral density.

6.2 Is the published theory completely positive?

The gravitational action is a positive weight on the scalar and tensor modes [11, 12, 13], but its vector weight $-\frac{\alpha}{4}p^2k^2|V|^2$ is negative on spacelike p, the strip itself. With matter added, Oppenheim and Sajjad find that the vector sector favours recoherence of the momentum density, so that "even if the gravitational vector sector is somehow fixed, the evolution would no longer be positive", and that "it is not fully clear to us how the vectorial sector is to be handled" [13]. They name remedies: (i) restricting the path integral to exclude tachyonic V, at the price of a non-Markovian noise kernel; (ii) a preferred frame, worked out as a unit timelike vector n^μ in the diffusion matrix, either in the vector and scalar sectors (their case A) or in the vector sector only (case C); (iii) an additional field; and (iv) restricting the kernel to timelike p [13]. The result does not depend on which, if any, is adopted. Define "the scalar sector stays Lorentz invariant" operationally: on the physical record space $D_2^{-1} = a_2(p^2)P^{(2)} + a_0(p^2)P^{(0s)} + \Delta$ with $\Delta \succeq 0$ vanishing, to $O(u/k^2)$, on the $B = 0$ source (14), and likewise for D_0. Then:

- *either* the matter-coupled theory is CP on the record space, which its authors leave open; its scalar sector is positive [12, 13], the strip noise is defined, and the theory is excluded as it stands, by $\geq 1.1 \times 10^8$ and by 2.5×10^3 in its constant-coupling window;
- *or* it needs a CP completion, and every completion whose scalar sector stays Lorentz invariant is excluded by the same margins.

The remedies fall on one side or the other. Remedy (i) and case C act on the vector sector only. With a positive vector sector the sectors are orthogonal, $\langle \Phi V \rangle = 0$, so the scalar floor on the $B = 0$ source is untouched and the scalar noise is a lower bound on the force noise; case C adds a non-negative term to the explicit channel. A non-Markovian kernel keeps the hypotheses of Lemma 2.1, which needs translation invariance and positivity on every past region, not Markovianity. Remedy (iv) sets the kernel to zero at spacelike p. If that is the decoherence kernel, Lemma 2.1 forces $D_1 = 0$ there, that is, no static Newtonian force. If it is the weight D_2^{-1}, the diffusion is infinite at spacelike p, as the authors note for V, and static Newton is lost in noise. If it is the noise kernel itself, the natural reading of a restriction on correlations, then $D_2 = 0$ at spacelike p and (6) forces $D_0 = \infty$ wherever $D_1 \neq 0$, which is also excluded. Static Newton, a drift with bounded noise at spacelike p, therefore excludes (iv) on the full kernel. Case A puts n^μ into the scalar weight and survives the scalar strip; it does not survive the tensor

Table 2. *Structures and the two strips. "Frame in the law" means a frame written into the dynamics; the other carriers are the classical record, the matter, and the state of an extra field. The tensor column assumes a tensor sector that is that of general relativity at keV momenta and assumptions (T1)–(T3).*

Structure	Scalar strip	Tensor strip
Conservation $k_\mu T^{\mu\nu} = 0$	response floor only	no
Constraint as a weight (covariant DeWitt)	no	no
Constraint as a frame law [7]	removed	no
Unit timelike vector in the scalar and vector weights (case A)	removed	no
Unit timelike vector in the vector weight only (case C)	no	no
Excluding tachyonic V, non-Markovian kernel	no	no
Kernel restricted to timelike p	excluded by static Newton	
Preferred foliation with a light-cone tensor	removed	no
Any p^2-only regulator or running	no	no
Finite age, cosmological friction	no	no
Perturbative or matter-induced scalar mass	no	no
Frame cut-off on k or ω below 1 keV in the full drift	removed	removed
Form factor in the matter four-velocity	removed	removed
Tensor dispersion off the cone below ~ 10 eV ($z = 3$)	—	removed
Matter coupling blind to the atom's stress	no	removed
Relational clock field [28]	open	open

strip (Section 5). Remedy (iii) is not specified; one realisation, a relational clock field [28], is open (Table 2).

6.3 Regulators, running couplings and the mean field

Every Lorentz-invariant regulator and every running of the couplings acts on u, and Theorem 3.2 holds for arbitrary $a_2(p^2)$, $a_0(p^2)$; the invariant ε and modulus-squared prescriptions of [12] are covered by rows 1 and 3 of Table 1. A cut-off on D_0 alone violates the trade-off unless the Newton coupling is cut as well; the floor holds wherever $|f|$ is tested, which is why the band is 1 mm–1 m. An invariant modification of the mean-field operator, $(1 + \ell^2\Box)G_{\mu\nu} = 8\pi G T_{\mu\nu}$, multiplies the product (18) by $(1 - \ell^2 u)^2$, which is one on the band for ℓ well below a millimetre.

6.4 A mass for the scalar

A mass acts on u, and the scalar floor becomes $\frac{2}{3}\int_1^{10^6} du/(u+m^2)$. At $m^2 = 10^6$ m^{-2} the margin is still 6.0×10^6; it reaches one at $m^2 = 8.6 \times 10^{12}$ m^{-2}, a Compton length of 0.34 μm. Masses induced by matter are far smaller: $8\pi G\rho/c^2 = 5.5 \times 10^{-23}$ m^{-2} in liquid xenon. The DeWitt family has no gauge-invariant quadratic mass term at linear order. A non-perturbative scalar mass above $(0.34\ \mu\text{m})^{-1}$ is the one Lorentz-invariant escape from the scalar strip within the Gaussian class at linear order. It does not touch the tensor floor.

6.5 What shields the scalar strip

Table 2 lists what removes each strip. Every structure that removes the scalar strip is a frame on the scalar sector near the light cone: in the law, carried by the record, carried by the matter's own four-momentum, or carried by the state of an extra field. A matter-frame form factor

$\exp(-[p^2 + 2(p \cdot U)^2]R_0^2/2)$, with U the four-velocity of the matter, depends on U only through the change of invariant mass at a vertex, so it acts only on inelastic vertices; the floor becomes $\frac{2}{3}|F|^2/u$ and needs only $R_0 \geq 1.68$ Å. A relational clock field [28] regulates the vacuum heating of a covariant master equation through $n_\mu \propto \partial_\mu T$, a frame carried by the state; whether the X-ray data bound its resolution is open.

6.6 Other inputs

- *Linearisation.* The strip has invariant lengths of at least 1 mm, deep infrared in a Lorentz-invariant theory. A non-linear effect that changes the scalar propagator on the strip without touching the static sector would shield; none is known.
- *Cosmological friction* acts only at $u \lesssim u_\eta = 1.0 \times 10^{-14}$ m^{-2} (140 keV), far below the band.
- *The atom.* An isolated local-density xenon atom with closure. For near-cone noise the force is not uniform on the atom ($qa \gg 1$), so A_{cone} rests on closure, which is least accurate below the xenon K edge; the E^2 template is dominated by 50–140 keV. Manti *et al.*'s isotropic form of the rate moves every margin by 1.4%.
- *The class.* The theorems are stated for Gaussian, continuous CQ dynamics. For jump-class dynamics [15, 27] the trade-off is a statement about second moments and holds as it stands [15]. For a marked Poisson record with no cross-correlation between the record and the explicit noise, the floors of Theorems 3.2 and 5.1 hold for the second cumulant. The charges that radiate in Eq. (21) respond linearly to the force, so the mean photon yield is the two-point value for any noise statistics. Non-Gaussian noise can change only how the photons are bunched. A jump-class theory therefore escapes only if all but a fraction of order the inverse margin of the floor's photons arrive in multi-photon bursts that the XENONnT selection rejects, at unchanged mean radiated power. We do not exclude such a theory here.

7 Discussion

7.1 Where this disagrees with the authors

Oppenheim and Sajjad conclude that covariant CQ gravity is viable given current bounds, with care needed over non-linearities, non-Markovianity and scale invariance [13]. Their assessment rests on the static and low-frequency scalar sector (matter-wave coherence, LISA Pathfinder) and on the tensor sector at gravitational-wave frequencies. They already find that constant couplings would be ruled out by the integrated energy of the stochastic gravitational waves, and they escape that bound through running couplings [13]. Their assessment does not evaluate the matter noise at keV. That channel excludes the constant-coupling window again, by 2.5×10^3 from the field noise alone. The sharper disagreement is with the escape: running couplings act on u, and the floor (row 1 of Table 1) holds for every running in p^2. The disagreement is one channel, and it is testable.

7.2 What it means for "is gravity classical?"

The result is narrower than the question. It does not show that gravity is quantized. It excludes Gaussian CQ dynamics in which Newton's force is carried by the classical field and either the scalar sector is Lorentz invariant up to boosts of 1.1×10^7, or the tensor sector is that of general relativity at keV momenta. Within that class the argument is short: positivity ties decoherence to diffusion; Lorentz invariance, or the tensor propagator itself, puts the diffusion next to the light cone at every X-ray energy; slow matter feels it through its rest mass or through the Maxwell stress of its own Coulomb field; xenon converts it to keV photons with an order-one atomic factor; XENONnT sees none.

A classical theory of gravity that survives both exclusions needs, at once,

1. a frame or a mass in the scalar sector near the light cone (Table 2), and
2. a tensor sector that leaves general relativity below about 10 eV (for a $z = 3$ Lifshitz dispersion), or a matter coupling that does not see the atom's stress,

or its noise must be non-Gaussian and bunch the keV emission into multi-photon bursts (Section 6). Frame-fixed Newtonian CQ models, whose least decoherence is the Diósi–Penrose rate [6], remain subject to the published Diósi–Penrose X-ray bounds [8, 10] and are not excluded by this analysis.

7.3 Open

- A tensor conversion for liquid xenon with a relativistic atom, which would turn the tensor margin from an estimate into a bound.
- A gauge-invariant treatment of matter couplings that omit the Coulomb field's stress.
- A non-perturbative scalar mass, outside the DeWitt family at linear order.
- The resolution of a relational clock field [28].
- Non-linear effects on the strip that leave the static sector alone.

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